Physics Letters B 802 (2020) 135254

Contents lists available at ScienceDirect

Physics Letters B

www.elsevier.com/locate/physletb

A theory for scotogenic dark matter stabilised by residual gauge symmetry

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A R T I C L E I N F O

Article history:
Received 19 September 2019
Received in revised form 16 January 2020
Accepted 23 January 2020
Available online 25 January 2020
Editor: A. Ringwald

A B S T R A C T

Dark matter stability can result from a residual matter-parity symmetry, following naturally from the spontaneous breaking of the gauge symmetry. Here we explore this idea in the context of the $SU(3)_c \otimes SU(3)_L \otimes U(1)_X \otimes U(1)_N$ electroweak extension of the standard model. The key feature of our new scotogenic dark matter theory is the use of a triplet scalar boson with anti-symmetric Yukawa couplings. This naturally implies that one of the light neutrinos is massless and, as a result, there is a lower bound for the $0\nu\beta\beta$ decay rate.

1. Introduction

In order to account for the existence of cosmological dark matter, we need new particles not present in the Standard Model (SM) of particle physics. Moreover, new symmetries capable of stabilising the corresponding candidate particle on cosmological scales are also required. Here we focus on the so-called Weakly Interacting Massive Particles, or WIMPs, as dark matter candidates. Within supersymmetric schemes, WIMP stability follows from having a conserved R-parity symmetry [1]. Our present construction does not rely on supersymmetry nor on the imposition of any *ad hoc* symmetry to stabilise dark matter. It is also a more complete theory setup, in the sense that it naturally generates neutrino masses as well. These arise radiatively, thanks to the exchange of new particles in the "dark" sector. The procedure is very well-motivated since neutrino masses are anyways necessary to account for neutrino oscillation data [2].

Here we follow an alternative approach that naturally incorporates neutrino mass right from the beginning. This is provided by scotogenic dark matter schemes. These are "low-scale" models of neutrino mass [3] where dark matter emerges as a radiative mediator of neutrino mass generation. In this case, the symmetry stabilising dark matter is also responsible for the radiative origin of neutrino masses in a very elegant way [4]. Yet, in this case too, a dark matter stabilisation symmetry is introduced in an *ad hoc* manner. The need for such "dark" symmetry is a generic feature also of other scotogenic schemes, such as the generalization proposed in [5,6].

Extending the $SU(3)_c \otimes SU(2)_L \otimes U(1)_Y$ gauge symmetry can provide a natural setting for a theory of dark matter where stabilisation can be automatic [7–9]. Such electroweak extensions involve the $SU(3)_L$ gauge symmetry, which also provides an "explanation" of the number of quark and lepton families from the anomaly cancellation requirement [10–12]. For recent papers using the $SU(3)_L$ gauge symmetry see Refs. [13–22]. These theories can also, in some cases, be made consistent with unification of the gauge couplings [23,24] and/or with the existence of left-right gauge symmetry [25,26]. In the extended electroweak gauge symmetry models discussed in [7,8] the stability of dark matter results from the presence of a matter-parity symmetry, M_P, a non-supersymmetric version of R-parity, that is a natural consequence of the spontaneous breaking of the extended gauge symmetry.

The purpose of this letter is to improve upon the proposal in [9] in two ways. First, we simplify the particle content. Compared with Ref. [9] no extra vector-like fermions nor scalar $SU(3)_L$ sextets are needed. Instead, the matter parity odd, third component (N_L) of the

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https://doi.org/10.1016/j.physletb.2020.135254

Table 1
Particle content of the theory. Here $a = 1, 2, 3$ and $i = 1, 2$ are family indices.

Field	$SU(3)_c$	$SU(3)_L$	$U(1)_X$	$U(1)_N$	Q	$B-L$	$M_P = (-1)^{3(B-L)+2s}$
q_{iL}	**3**	$\bar{\mathbf{3}}$	0	0	$(-\frac{1}{3}, \frac{2}{3}, -\frac{1}{3})^T$	$(\frac{1}{3}, \frac{1}{3}, -\frac{2}{3})^T$	$(++-)^T$
q_{3L}	**3**	**3**	$\frac{1}{3}$	$\frac{2}{3}$	$(\frac{2}{3}, -\frac{1}{3}, \frac{2}{3})^T$	$(\frac{1}{3}, \frac{1}{3}, \frac{4}{3})^T$	$(++-)^T$
u_{aR}	**3**	**1**	$\frac{2}{3}$	$\frac{1}{3}$	$\frac{2}{3}$	$\frac{1}{3}$	$+$
d_{aR}	**3**	**1**	$-\frac{1}{3}$	$\frac{1}{3}$	$-\frac{1}{3}$	$\frac{1}{3}$	$+$
U_{3R}	**3**	**1**	$\frac{2}{3}$	$\frac{4}{3}$	$\frac{2}{3}$	$\frac{4}{3}$	$-$
D_{iR}	**3**	**1**	$-\frac{1}{3}$	$-\frac{2}{3}$	$-\frac{1}{3}$	$-\frac{2}{3}$	$-$
l_{aL}	**1**	**3**	$-\frac{1}{3}$	$-\frac{2}{3}$	$(0, -1, 0)^T$	$(-1, -1, 0)^T$	$(++-)^T$
e_{aR}	**1**	**1**	-1	-1	-1	-1	$+$
ν_{iR}	**1**	**1**	0	-4	0	-4	$-$
ν_{3R}	**1**	**1**	0	5	0	5	$+$
N_{aR}	**1**	**1**	0	0	0	0	$-$
η	**1**	**3**	$-\frac{1}{3}$	$\frac{1}{3}$	$(0, -1, 0)^T$	$(0, 0, 1)^T$	$(++-)^T$
ρ	**1**	**3**	$\frac{2}{3}$	$\frac{1}{3}$	$(1, 0, 1)^T$	$(0, 0, 1)^T$	$(++-)^T$
χ	**1**	**3**	$-\frac{1}{3}$	$-\frac{2}{3}$	$(0, -1, 0)^T$	$(-1, -1, 0)^T$	$(--+)^T$
σ	**1**	**1**	0	2	0	2	$+$
ζ	**1**	**3**	$\frac{2}{3}$	$\frac{7}{3}$	$(1, 0, 1)^T$	$(2, 2, 3)^T$	$(+, +, -)^T$
ξ	**1**	**3**	$\frac{2}{3}$	$\frac{4}{3}$	$(1, 0, 1)^T$	$(1, 1, 2)^T$	$(-, -, +)^T$

$SU(3)_L$ lepton triplet plays the role of dark fermion with its Dirac partner being a new $SU(3)_L$ singlet N_R. Moreover, the dark sextet scalar particles are replaced by an $SU(3)_L$ scalar triplet. As seen in Fig. 1 these particles are enough to implement the scotogenic scenario. Moreover, in contrast to the proposal in Ref. [9], here we predict that one of the light neutrinos is massless. This feature arises in a novel way when compared to other schemes in the literature. So far most realistic theories where one of the neutrinos is (nearly) massless typically involve "missing partner" schemes, such as the "incomplete" seesaw mechanism [27] or similar radiative mechanisms [28].

The paper is organized as follows. In Sec. 2 we present the model, while the loop-induced neutrino masses are discussed in Sec 3. The symmetry breaking sector, scalar potential and mass spectrum are discussed in Sec. 4. Concerning phenomenology, in Secs. 5 and 6, we briefly comment on dark matter and the predicted lower bound for the $0\nu\beta\beta$ decay rate, as well as FCNC and collider signatures. Finally, our conclusions are presented in Sec. 7.

2. Our model

Here we give the main features of the model, based on the $SU(3)_c \otimes SU(3)_L \otimes U(1)_X \otimes U(1)_N$ gauge invariance. The electric charge and $B-L$ generators are given by

$$Q = T_3 - \frac{1}{\sqrt{3}} T_8 + X \,, \tag{1}$$

$$B - L = -\frac{2}{\sqrt{3}} T_8 + N \,, \tag{2}$$

where T_m, with $m = 1, 2, 3, ..., 8$, are the $SU(3)_L$ generators, whereas X and N are the $U(1)_X$ and $U(1)_N$ generators, respectively. Notice that, due to the extra $U(1)_N$ symmetry, the $B-L$ symmetry is fully gauged. The SM $SU(2)_L$ doublet quarks and leptons reside inside the $SU(3)_L$ anti-triplet q_{iL}, triplet q_{3L}; $i = 1, 2$ and l_{aL}; $a = 1, 2, 3$ and their field decomposition is given by:

$$q_{iL} = \begin{pmatrix} d_{iL} \\ -u_{iL} \\ D_{iL} \end{pmatrix}, \qquad q_{3L} = \begin{pmatrix} u_{3L} \\ d_{3L} \\ U_{3L} \end{pmatrix}, \qquad l_{aL} = \begin{pmatrix} \nu_{aL} \\ e_{aL} \\ N_{aL} \end{pmatrix}, \tag{3}$$

whereas their $SU(2)_L$ singlet partners are given by u_{aR}, d_{aR} and e_{aR} respectively. The full particle content of the model along with the corresponding charges is summarised in Table 1.

Symmetry breaking takes place through the non-vanishing vacuum expectation values (vevs) as given below,

$$\langle\sigma\rangle = \frac{v_\sigma}{\sqrt{2}}, \langle\chi\rangle = \frac{1}{\sqrt{2}}(0, 0, w)^T, \tag{4}$$
$$\langle\eta\rangle = \frac{1}{\sqrt{2}}(v_1, 0, 0)^T, \langle\rho\rangle = \frac{1}{\sqrt{2}}(0, v_2, 0)^T, \langle\zeta\rangle = \frac{1}{\sqrt{2}}(0, v_2', 0)^T .$$

The top vevs break $SU(3) \otimes SU(3)_L \otimes U(1)_X \otimes U(1)_N$ to the SM gauge symmetry with $v_\sigma, w \gg v_{EW}$, while $v_{EW} = (v_1^2 + v_2^2 + v_2'^2)^{1/2} = 246$ GeV leads to electroweak breaking. Note that, while w breaks $SU(3)_L \otimes U(1)_X$, v_σ breaks $U(1)_N$. When σ and χ acquire similar vevs, $v_\sigma \sim w$, the two steps of the symmetry breaking process occur at the same time.

$$SU(3)_c \otimes SU(3)_L \otimes U(1)_X \otimes U(1)_N \xrightarrow{v_\sigma, w} SU(3)_c \otimes SU(2)_L \otimes U(1)_Y \otimes M_P \,. \tag{5}$$

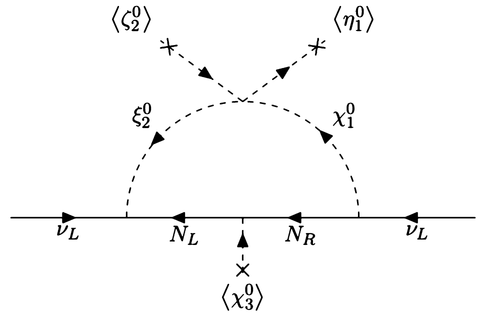

Fig. 1. One-loop "scotogenic" neutrino mass.

The last step takes place when the first and second components of the triplets acquire vevs, and we are left with

$$SU(3)_c \otimes SU(2)_L \otimes U(1)_Y \otimes M_P \xrightarrow{v_1, v_2, v_2'} SU(3)_c \otimes U(1)_Q \otimes M_P \,, \tag{6}$$

in such a way that $(v_1^2 + v_2^2 + v_2'^2)^{1/2} = v_{EW}$, the electroweak scale.

This process leaves at the end a matter-parity symmetry, M_P, defined as

$$M_P = (-1)^{3(B-L)+2s} \,. \tag{7}$$

Notice the important fact that only the M_P-even scalar fields get vevs. This implies that matter-parity remains as an absolutely conserved residual gauge symmetry even after spontaneous symmetry breaking, implying that the lightest amongst the M_P-odd particles is stable. Here we notice that the presence of the nonvanishing vev v_σ breaks $U(1)_N$ at a potentially large scale, preventing the appearance of a light Z' gauge boson.

3. Neutrino masses

Taking into account the leptons and scalars present in our model, as shown in Table 1, the following Yukawa sector can be written down

$$-\mathcal{L}_{lep} = y^e_{ab}\overline{l_{aL}}\rho e_{bR} + y^N_{ab}\overline{l_{aL}}\chi N_{bR} + h_{ab}\overline{l_{aL}}(l_{bL})^c\xi^* + \frac{(m_N)_{ab}}{2}\overline{(N_{aR})^c}N_{bR} + h.c. \,, \tag{8}$$

where y^e, y^N, h and m_N are complex 3×3 matrices, where m_N is symmetric, due to the Pauli principle. In contrast, due again to the symmetry structure of the theory, the Yukawa coupling matrix h is anti-symmetric in family space. Notice that this anti-symmetric Yukawa coupling was first proposed in [29]. While the original scheme is no longer viable, given the current neutrino oscillation data, the new construction provides a consistent variant that also accounts for WIMP dark matter in a scotogenic way, i.e. dark matter emerges as a neutrino mass mediator, see Fig. 1.

Notice that, while the fields ν_{iR} and ν_{3R} with non-standard charges [30–32] are necessary in order to ensure anomaly cancellation, such choice of charges forbids their coupling to the other leptons as well as scalars, justifying their absence from the Lagrangian given above.[1]

The first term in Eq. (8) generates a mass term to the charged leptons when ρ acquires a vev:

$$M^e = y^e \frac{v_2}{\sqrt{2}} \,, \tag{9}$$

where the family indices have been omitted. The neutral leptons N_{iL} and N_{iR} mass matrix is given as

$$M^N = \frac{1}{2}\begin{pmatrix} 0 & y^N w \\ (y^N)^T w & m_N \end{pmatrix} \tag{10}$$

in the basis $(N_L, (N_R)^c)^T$. Such a matrix is diagonalised (in the one family approximation) as

$$\begin{pmatrix} N_1 \\ N_2 \end{pmatrix} = \begin{pmatrix} \cos\theta_N & -\sin\theta_N \\ \sin\theta_N & \cos\theta_N \end{pmatrix}\begin{pmatrix} N_L \\ (N_R)^c \end{pmatrix} \text{with} \tan(2\theta_N) = \frac{2y^N w}{m_N}. \tag{11}$$

Turning to the light neutrinos ν_L, it is easy to see that tree-level mass terms could be generated if the first component of χ or the second of ξ acquired a vev. This, however, does not occur as a result of the assumed pattern of vevs and this, in turn, is dynamically consistent with the minimization of the potential. This way matter-parity conservation emerges as a residual symmetry.

Neutrino masses are radiatively generated by the one-loop diagram in Fig. 1. The relevant scalar interaction is the one governed by the λ_1, see Eq. (16), and leads to

[1] ν_{iR} and ν_{3R} masses could be generated, for example, by coupling them to new scalars transforming as $(\mathbf{1}, \mathbf{1}, 0, 8)$ and $(\mathbf{1}, \mathbf{1}, 0, -10)$, respectively. Note that matter-parity conservation would not be spoiled when such scalars acquire vevs.

$$m_\nu^{ab} = \frac{1}{8\pi^2} h^{*ac} s_N c_N c_1 \left\{ m_{N_1} \left[s_{S_2} c_{S_2} \left(Z\left(\frac{m_{S_1}^2}{m_{N_1}^2}\right) - Z\left(\frac{m_{S_2}^2}{m_{N_1}^2}\right)\right) - s_{A_2} c_{A_2} \left(Z\left(\frac{m_{A_1}^2}{m_{N_1}^2}\right) - Z\left(\frac{m_{A_2}^2}{m_{N_1}^2}\right)\right)\right] \right.$$
$$\left. - m_{N_2} \left[s_{S_2} c_{S_2} \left(Z\left(\frac{m_{S_1}^2}{m_{N_2}^2}\right) - Z\left(\frac{m_{S_2}^2}{m_{N_2}^2}\right)\right) - s_{A_2} c_{A_2} \left(Z\left(\frac{m_{A_1}^2}{m_{N_2}^2}\right) - Z\left(\frac{m_{A_2}^2}{m_{N_2}^2}\right)\right)\right] \right\}_{cd} y^{N*db} + \{a \leftrightarrow b\}, \tag{12}$$

where $s_x \equiv \sin\theta_x$, $c_x \equiv \cos\theta_x$, and the loop function $Z(x)$ is defined as

$$Z(x) = \frac{x}{1-x} \ln x. \tag{13}$$

An important point to note here is that owing to the antisymmetry of the Yukawa matrix h, the resulting neutrino mass matrix of Eq. (12) is of rank two. This implies that when rotated to the mass basis, only two neutrinos acquire mass and one remains massless. This unique feature provides a novel origin for the masslessness of one neutrino[2] that should be contrasted with the usual models for one massless neutrino, which typically rely on missing partner mechanisms.

Note that the matter-parity odd neutral fermions $\left(N_L, (N_R)^c\right)$ are obtained from Eq. (11), while the scalar masses $m_{1,2}$ are given in Eq. (23) and the mixing angles of $\left(\xi_2^0, \chi_1^0, \eta_3^0\right)_{S,A}$ come from Eq. (24).

It is worth pointing out that, in addition to the usual loop suppression characteristic of scotogenic models, our result in Eq. (12) is further suppressed by the factor $c_1 \sim v_1/w \ll 1$, see Eq. (20). This is needed in order to identify the physical mass eigenstates associated with the scalar mediators in the scotogenic loop.

All fields running inside the neutrino mass loop are odd under matter-parity. The exact conservation of this symmetry implies that the lightest among the M_P-odd particles is stable, and therefore can play the role of dark matter. Thus, the present model generates "scotogenic" neutrino masses, with the crucial dark matter stabilising symmetry emerging naturally as a residual subgroup of the original gauge symmetry.

4. Scalar sector

In addition to the three $SU(3)_L$ triplets η, χ, ρ our model employs two others, ξ and ζ. The scalar triplets can be decomposed into

$$\eta, \chi \equiv \begin{pmatrix} \eta_1^0 \\ \eta_2^- \\ \eta_3^0 \end{pmatrix}, \begin{pmatrix} \chi_1^0 \\ \chi_2^- \\ \chi_3^0 \end{pmatrix} \rho, \xi, \zeta \equiv \begin{pmatrix} \rho_1^+ \\ \rho_2^0 \\ \rho_3^+ \end{pmatrix}, \begin{pmatrix} \xi_1^+ \\ \xi_2^0 \\ \xi_3^+ \end{pmatrix}, \begin{pmatrix} \zeta_1^+ \\ \zeta_2^0 \\ \zeta_3^+ \end{pmatrix}, \tag{14}$$

and the neutral components, as well as the scalar singlet σ, can be further decomposed into their CP-even (S) and CP-odd (A) parts, in such a way that for a given neutral scalar field s_i^0, we have

$$s_i^0 \equiv \frac{1}{\sqrt{2}}(v_{s_i} + S_{s_i} + iA_{s_i}), \tag{15}$$

with s denoting generically all the scalars, and $v_{\eta_1} = v_1$, $v_{\rho_2} = v_2$, $v_{\chi_3} = w$, $v_{\zeta_2} = v_2'$, and $v_\sigma = v_\sigma$, whereas all the other vevs vanish, as already discussed in Eq. (4).

Given the five scalar triplets and the singlet in Table 1, the scalar potential can be written as

$$V = \sum_s \left[\mu_s^2 (s^\dagger s) + \frac{\lambda_s}{2}(s^\dagger s)^2\right] + \sum_{s_1,s_2}^{s_1>s_2} \left[\lambda_{s_1 s_2}(s_1^\dagger s_1)(s_2^\dagger s_2)\right] + \sum_{t_1,t_2}^{t_1>t_2} \left[\lambda'_{t_1 t_2}(t_1^\dagger t_2)(t_2^\dagger t_1)\right]$$
$$+ \frac{\mu_1}{\sqrt{2}} \eta\rho\chi + \frac{\mu_2}{\sqrt{2}}(\zeta^\dagger\rho)\sigma + \lambda_1(\chi^\dagger\eta)(\zeta^\dagger\xi) + \lambda_2(\chi^\dagger\xi)(\zeta^\dagger\eta) + \lambda_3(\chi^\dagger\eta)(\xi^\dagger\rho) + \lambda_4(\chi^\dagger\rho)(\xi^\dagger\eta) + \lambda_5(\eta\zeta\chi)\sigma^* + h.c., \tag{16}$$

where t only varies through all the scalar triplets: $t = \eta, \chi, \rho, \xi, \zeta$, while s varies through all the scalars, *i.e.* the triplets in t plus the singlet σ.

By minimising the scalar potential, we obtain the "tadpole" conditions

$$\mu_1 v_2 w + \lambda_5 v_2' w v_\sigma + v_1 \left(2\mu_\eta^2 + \lambda_\eta v_1^2 + \lambda_{\eta\rho} v_2^2 + \lambda_{\eta\zeta} v_2'^2 + \lambda_{\eta\chi} w^2 + \lambda_{\eta\sigma} v_\sigma^2\right) = 0, \tag{17}$$
$$\mu_1 v_1 w + \mu_2 v_2' v_\sigma + v_2 \left[2\mu_\rho^2 + \lambda_{\eta\rho} v_1^2 + \lambda_\rho v_2^2 + (\lambda_{\rho\zeta} + \lambda_{\rho\zeta 2}) v_2'^2 + \lambda_{\rho\chi} w^2 + \lambda_{\rho\sigma} v_\sigma^2\right] = 0,$$
$$\mu_1 v_1 v_2 + \lambda_5 v_1 v_2' v_\sigma + w \left(2\mu_\chi^2 + \lambda_{\eta\chi} v_1^2 + \lambda_{\rho\chi} v_2^2 + \lambda_{\chi\zeta} v_2'^2 + \lambda_\chi w^2 + \lambda_{\chi\sigma} v_\sigma^2\right) = 0,$$
$$\mu_2 v_2 v_\sigma + \lambda_5 v_1 v_\sigma w + v_2' \left[2\mu_\zeta^2 + \lambda_{\eta\zeta} v_1^2 + (\lambda_{\rho\zeta} + \lambda_{\rho\zeta 2}) v_2^2 + \lambda_\zeta v_2'^2 + \lambda_{\chi\zeta} w^2 + \lambda_{\zeta\sigma} v_\sigma^2\right] = 0,$$
$$\mu_2 v_2 v_2' + \lambda_5 v_1 v_2' w + v_\sigma \left(2\mu_\sigma^2 + \lambda_{\eta\sigma} v_1^2 + \lambda_{\rho\sigma} v_2^2 + \lambda_{\zeta\sigma} v_2'^2 + \lambda_{\chi\sigma} w^2 + \lambda_\sigma v_\sigma^2\right) = 0,$$

[2] This is reminiscent of the proposal in [29], currently ruled out by the oscillation data.

through which $\mu_\eta, \mu_\rho, \mu_\chi, \mu_\zeta$ and μ_σ can be eliminated from the potential. Nine out of the initial degrees of freedom in the scalar sector are absorbed as longitudinal components of the massive gauge vector bosons, $Z, Z', Z'', U^0, (U^0)^\dagger, W^\pm, V^\pm$. The remaining scalar fields become massive, as we now discuss.

First, we focus on the scalar fields that enter the neutrino mass loop, for which we show the corresponding mass matrices and diagonalise them in Sec. 4.1, providing the mass eigenvalues and eigenstates. For the other scalars, the mass matrices are given in Sec. 4.2.

4.1. Neutrino-mass-mediator scalars

The scalar fields relevant to the neutrino mass loop in Fig. 1 are part of the set of the M_P-odd neutral fields and can be grouped together into a CP-even and a CP-odd set: $(S_{\xi_2}, S_{\chi_1}, S_{\eta_3})$ and $(A_{\xi_2}, A_{\chi_1}, A_{\eta_3})$, respectively. In such bases, we can write down the following squared mass matrices

$$M^2_{S,A} = \frac{1}{2}\begin{pmatrix} a_{11} & a_{12} & a_{13} \\ a_{12} & a_{22} & a_{23} \\ a_{13} & a_{23} & a_{33} \end{pmatrix}_{S,A} , \tag{18}$$

where the elements a_{ij} are defined as

$$\begin{aligned}
(a_{11})_{S,A} &= a_{11} = \lambda_{\eta\xi} v_1^2 + (\lambda_{\rho\xi} + \lambda_{\rho\xi 2}) v_2^2 + (\lambda_{\xi\zeta} + \lambda_{\xi\zeta 2}) v_2'^2 + \lambda_{\chi\xi} w^2 + \lambda_{\xi\sigma} v_\sigma^2 + 2\mu_\xi^2 , \\
(a_{22})_{S,A} &= a_{22} = \lambda_{\eta\chi 2} v_1^2 - \frac{v_1}{w}(\lambda_5 v_2' v_\sigma + \mu_1 v_2), \\
(a_{33})_{S,A} &= a_{33} = \lambda_{\eta\chi 2} w^2 - \frac{w}{v_1}(\lambda_5 v_2' v_\sigma + \mu_1 v_2), \\
(a_{12})_{S,A} &= v_1(\lambda_1 v_2' \pm \lambda_3 v_2) , \\
(a_{13})_{S,A} &= w(\lambda_3 v_2 \pm \lambda_1 v_2') , \\
(a_{23})_{S,A} &= \pm(\lambda_{\eta\chi 2} v_1 w - \mu_1 v_2 - \lambda_5 v_2' v_\sigma) .
\end{aligned} \tag{19}$$

Each matrix has a vanishing eigenvalue associated with a would-be Goldstone boson that is absorbed by the gauge sector, more specifically by the complex neutral gauge field U_0. We can find the massless eigenstate by rotating the second and third components of both the CP-even and CP-odd basis by

$$(\theta_1)_{S,A} = \pm \arctan\left(\frac{w}{v_1}\right) , \tag{20}$$

respectively. After these transformations, the matrices in Eq. (18) become

$$\tilde{M}^2_{S,A} = \frac{1}{2}\begin{pmatrix} a_{11} & x a_{12} & 0 \\ x a_{12} & x^2 a_{22} & 0 \\ 0 & 0 & 0 \end{pmatrix}_{S,A} , \text{with} x^2 = \frac{v_1^2 + w^2}{v_1^2} . \tag{21}$$

Such matrices can be finally diagonalised by rotating the two first components of each basis by

$$(\theta_2)_{S,A} = \frac{1}{2}\arctan\left[\frac{2x(a_{12})_{S,A}}{a_{11} - x^2 a_{22}}\right] , \tag{22}$$

respectively. By doing so, we obtain the eigenvalues

$$(m^2_{1,2})_{S,A} = \frac{1}{2}\left[a_{11} + x^2 a_{22} \pm \sqrt{(a_{11} - x^2 a_{22})^2 + 4x^2 (a_{12})^2_{S,A}}\right], (m_3^2)_{S,A} = 0. \tag{23}$$

In summary, the mass and flavour states can be related as

$$\begin{pmatrix} (S,A)_{m_1} \\ (S,A)_{m_2} \\ (S,A)_{m_3} \end{pmatrix} = \begin{pmatrix} \cos\theta_2 & \sin\theta_2 & 0 \\ -\sin\theta_2 & \cos\theta_2 & 0 \\ 0 & 0 & 1 \end{pmatrix}_{S,A} \begin{pmatrix} 1 & 0 & 0 \\ 0 & \cos\theta_1 & \sin\theta_1 \\ 0 & -\sin\theta_1 & \cos\theta_1 \end{pmatrix}_{S,A} \begin{pmatrix} (S,A)_{\xi_2} \\ (S,A)_{\chi_1} \\ (S,A)_{\eta_3} \end{pmatrix} . \tag{24}$$

4.2. Mass matrices of the other scalars

In this section we present the squared mass matrices associated with the scalar fields that do not take part in the neutrino mass loop. The CP-even and M_P-even neutral fields can be grouped in the basis $(S_{\eta_1}, S_{\rho_2}, S_{\chi_3}, S_{\zeta_2}, S_\sigma)$, so that we have the following symmetric squared mass matrix

$$M^2_{S_2} = \frac{1}{2}\begin{pmatrix} b_{11} & 2\lambda_{\eta\rho} v_1 v_2 + \mu_1 w & 2\lambda_{\eta\chi} v_1 w + \mu_1 v_2 + \lambda_5 v_2' v_\phi & 2\lambda_{\eta\zeta} v_1 v_2' + \lambda_5 v_\sigma w & 2\lambda_{\eta\sigma} v_1 v_\sigma + \lambda_5 v_2' w \\ \star & b_{22} & 2\lambda_{\rho\chi} v_2 w + \mu_1 v_1 & 2(\lambda_{\rho\zeta} + \lambda_{\rho\zeta 2}) v_2 v_2' + \mu_2 v_\sigma & 2\lambda_{\rho\sigma} v_2 v_\sigma + \mu_2 v_2' \\ \star & \star & b_{33} & 2\lambda_{\chi\zeta} v_2' w + \lambda_5 v_1 v_\sigma & 2\lambda_{\chi\sigma} v_\sigma w + \lambda_5 v_1 v_2' \\ \star & \star & \star & b_{44} & 2\lambda_{\zeta\sigma} v_2' v_\sigma + \mu_2 v_2 + \lambda_5 v_1 w \\ \star & \star & \star & \star & b_{55} \end{pmatrix}, \tag{25}$$

with the diagonal elements given by

$$\begin{aligned}
b_{11} &= 2\lambda_\eta v_1^2 - \frac{w}{v_1}\left(\mu_1 v_2 + \lambda_5 v_2' v_\sigma\right) , \\
b_{22} &= 2\lambda_\rho v_2^2 - \frac{\mu_1 v_1 w + \mu_2 v_2' v_\sigma}{v_2} , \\
b_{33} &= 2\lambda_\chi w^2 - \frac{v_1}{w}\left(\mu_1 v_2 + \lambda_5 v_2' v_\sigma\right) , \\
b_{44} &= 2\lambda_\zeta v_2'^2 - \frac{v_\sigma}{v_2'}\left(\mu_2 v_2 + \lambda_5 v_1 w\right) , \\
b_{55} &= 2\lambda_\sigma v_\sigma^2 - \frac{v_2'}{v_\sigma}\left(\mu_2 v_2 + \lambda_5 v_1 w\right) .
\end{aligned} \tag{26}$$

Upon diagonalisation, five non-vanishing masses appear associated with five physical scalars, one of which is the 125 GeV Higgs boson discovered at the LHC.

Taking into account now the CP-odd, M_P-even fields we obtain the squared mass matrix below, expressed in the basis $(A_{\eta_1}, A_{\rho_2}, A_{\chi_3}, A_{\zeta_2}, A_\sigma)$,

$$M_{A_2}^2 = \frac{1}{2}\begin{pmatrix}
-\frac{w}{v_1}(\mu_1 v_2 + \lambda_5 v_2' v_\sigma) & -\mu_1 w & -\mu_1 v_2 - \lambda_5 v_2' v_\sigma & -\lambda_5 v_\sigma w & \lambda_5 v_2' w \\
\star & -\frac{\mu_1 v_1 w + \mu_2 v_2' v_\sigma}{v_2} & -\mu_1 v_1 & \mu_2 v_\sigma & -\mu_2 v_2' \\
\star & \star & -\frac{v_1}{w}(\mu_1 v_2 + \lambda_5 v_2' v_\sigma) & -\lambda_5 v_1 v_\sigma & \lambda_5 v_1 v_2' \\
\star & \star & \star & -\frac{v_\sigma}{v_2'}(\mu_2 v_2 + \lambda_5 v_1 w) & \mu_2 v_2 + \lambda_5 v_1 w \\
\star & \star & \star & \star & -\frac{v_2'}{v_\sigma}(\mu_2 v_2 + \lambda_5 v_1 w)
\end{pmatrix} . \tag{27}$$

Three states remain massless and are absorbed by the neutral gauge bosons Z, Z', Z''. The other two states give rise to two massive CP-odd scalars.

At last, we consider the charged scalar fields. In the M_P-even basis $(\eta_2^\pm, \rho_1^\pm, \zeta_1^\pm, \xi_3^\pm)$, we can write the first squared mass matrix as

$$(M_1^\pm)^2 = \frac{1}{2}\begin{pmatrix}
c_{11} & \lambda_{\eta\rho 2} v_1 v_2 - \mu_1 w & \lambda_{\eta\zeta 2} v_1 v_2' - \lambda_5 v_\sigma w & \lambda_2 v_2' w \\
\lambda_{\eta\rho 2} v_1 v_2 - \mu_1 w & c_{22} & \lambda_{\rho\zeta 2} v_2 v_2' + \mu_2 v_\sigma & 0 \\
\lambda_{\eta\zeta 2} v_1 v_2' - \lambda_5 v_\sigma w & \lambda_{\rho\zeta 2} v_2 v_2' + \mu_2 v_\sigma & c_{33} & \lambda_2 v_1 w \\
\lambda_2 v_2' w & 0 & \lambda_2 v_1 w & c_{44}
\end{pmatrix} \tag{28}$$

with

$$\begin{aligned}
c_{11} &= \lambda_{\eta\rho 2} v_2^2 + \lambda_{\eta\zeta 2} v_2'^2 - \frac{w}{v_1}(\mu_1 v_2 + \lambda_5 v_2' v_\sigma) , \\
c_{22} &= \lambda_{\eta\rho 2} v_1^2 - \lambda_{\rho\zeta 2} v_2'^2 - \frac{\mu_1 v_1 w + \mu_2 v_2' v_\sigma}{v_2} , \\
c_{33} &= \lambda_{\eta\zeta 2} v_1^2 - \lambda_{\rho\zeta 2} v_2^2 - \frac{\mu_2 v_2 v_\sigma + \lambda_5 v_1 v_\sigma w}{v_2'} , \\
c_{44} &= \lambda_{\eta\xi} v_1^2 + \lambda_{\rho\xi} v_2^2 + \lambda_{\xi\zeta} v_2'^2 + (\lambda_{\chi\xi} + \lambda_{\chi\xi 2}) w^2 + \lambda_{\xi\sigma} v_\sigma^2 + 2\mu_\xi^2 .
\end{aligned} \tag{29}$$

Whereas in the M_P-odd basis $(\xi_1^\pm, \chi_2^\pm, \rho_3^\pm, \zeta_3^\pm)$, we write down

$$(M_2^\pm)^2 = \frac{1}{2}\begin{pmatrix}
d_{11} & \lambda_4 v_1 v_2 & \lambda_4 v_1 w & 0 \\
\lambda_4 v_1 v_2 & d_{22} & \lambda_{\rho\chi 2} v_2 w - \mu_1 v_1 & \lambda_{\chi\zeta 2} v_2' w - \lambda_5 v_1 v_\sigma \\
\lambda_4 v_1 w & \lambda_{\rho\chi 2} v_2 w - \mu_1 v_1 & d_{33} & \lambda_{\rho\zeta 2} v_2 v_2' + \mu_2 v_\sigma \\
0 & \lambda_{\chi\zeta 2} v_2' w - \lambda_5 v_1 v_\sigma & \lambda_{\rho\zeta 2} v_2 v_2' + \mu_2 v_\sigma & d_{44}
\end{pmatrix} \tag{30}$$

with

$$\begin{aligned}
d_{11} &= (\lambda_{\eta\xi} + \lambda_{\eta\xi 2}) v_1^2 + \lambda_{\rho\xi} v_2^2 + \lambda_{\xi\zeta} v_2'^2 + \lambda_{\chi\xi} w^2 + \lambda_{\xi\sigma} v_\sigma^2 + 2\mu_\xi^2 , \\
d_{22} &= \lambda_{\rho\chi 2} v_2^2 + \lambda_{\chi\zeta 2} v_2'^2 - \frac{v_1}{w}(\mu_1 v_2 + \lambda_5 v_2' v_\sigma), \\
d_{33} &= \lambda_{\rho\chi 2} w^2 - \lambda_{\rho\zeta 2} v_2'^2 - \frac{\mu_1 v_1 w + \mu_2 v_2' v_\sigma}{v_2}, \\
d_{44} &= \lambda_{\chi\zeta 2} w^2 - \lambda_{\rho\zeta 2} v_2^2 - \frac{v_\sigma}{v_2'}(\mu_2 v_2 + \lambda_5 v_1 w).
\end{aligned} \tag{31}$$

Each of the squared mass matrices above has a vanishing eigenvalue associated with a would-be Goldstone boson that will be absorbed by the charged gauge bosons $W^\pm$ and $V^\pm$. Finally, we are left with six heavy charged scalar fields in the model.

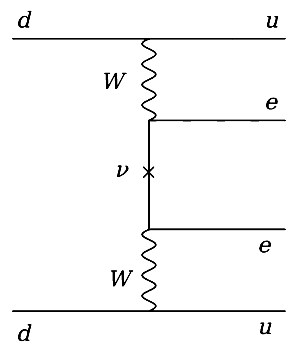

Fig. 2. Standard "mass-mechanism" $0\nu\beta\beta$ contribution.

5. Scotogenic dark matter neutrinoless double beta decay

This model can harbour a WIMP dark matter candidate that can be either scalar or a fermion.

Whatever the dark matter profile will be, we note the presence of new interactions, in addition to the standard processes of the simplest scotogenic model, which are primarily responsible for setting the dark matter relic abundance. In particular there are new t-channel processes involving the matter-parity-odd electrically neutral gauge boson connecting the same-charge components of the fermion triplets.

Concerning dark matter detection, let us recall that in the simplest scotogenic scenario [4], it proceeds primarily through the Higgs portal. In our model this portal has additional contributions, thanks to the presence of new scalar bosons. Moreover, since the dark sector particles carry $B-L$ charges, the usual Higgs portal is accompanied by a Z' portal. Which of the two portals will be dominant depends on the $B-L$ breaking scale, as well as on the various coupling strengths, particularly the Higgs-dark matter quartic coupling and $B-L$ gauge coupling strength g'. In the limit of large v_σ and w we recover the standard scotogenic dark matter phenomenology, which has been investigated before in [9]. In contrast, if the $B-L$ breaking scale is low the Z' may become significant. However, both scenarios have been well studied in the literature see, e.g. [32] for the Z' portal, so they will not be analysed here. Instead, we move directly to neutrinoless double beta decay, which presents interesting characteristic features.

Whether neutrinos are Majorana or Dirac fermions is still an open question. The case for Majorana neutrinos, as predicted by our model, can be undoubtedly established if neutrinoless double beta decay is ever observed [33]. The standard light neutrino-mediated $0\nu\beta\beta$ decay contribution is shown in Fig. 2. Its amplitude involves the lightest charged gauge boson $W^\pm$ exchange and hence is expressed in terms of the Fermi constant G_F, the typical momentum exchange p characterizing the process, and the effective Majorana mass $\langle m_{\beta\beta}\rangle$

$$\langle m_{\beta\beta}\rangle = |\cos^2\theta_{12}\cos^2\theta_{13}m_1 + \sin^2\theta_{12}\cos^2\theta_{13}m_2 e^{2i\phi_{12}} + \sin^2\theta_{13}m_3 e^{2i\phi_{13}}|\,, \tag{32}$$

is neatly expressed in the symmetric parametrization of the lepton mixing matrix [27] in terms of the mixing angles θ_{12} and θ_{13}, the physical Majorana phases [34] ϕ_{12} and ϕ_{13}, and the neutrino mass eigenvalues m_a obtained from Fig. 1.

It is well-known that, in a generic model, this amplitude can vanish as a result of destructive interference amongst the three light neutrinos. This actually can happen for normal-ordered neutrinos, currently preferred by oscillations [2]. An important feature that emerges from the structure of our model is that one of the light neutrinos is predicted to be massless. In this case, with a massless neutrino in the spectrum, $\langle m_{\beta\beta}\rangle$ is given in terms of just one free parameter, the relative Majorana phase: $\phi \equiv \phi_{12} - \phi_{13}$, all other parameters are fairly well-determined by the oscillation experiments. One can easily verify that in this case the effective Majorana mass $\langle m_{\beta\beta}\rangle$ never vanishes, even for the case of normal mass ordering, as shown in Fig. 3. Thus, thanks to the presence of a massless neutrino, our model is testable, at least for the inverted ordering (IO) case, which falls within the expected sensitivity of the upcoming next generation $0\nu\beta\beta$ decay experiments.

The top four horizontal bands in Fig. 3 represent the current experimental limits coming from CUORE ($\langle m_{\beta\beta}\rangle < 110-520$ meV) [35], EXO 200 Phase II ($93-286$ meV) [36], Gerda Phase II ($120-260$ meV) [37] and Kamland Zen ($61-165$ meV) [38]. The widths of these bands reflect uncertainties in nuclear matrix elements. The lower bands show the future sensitives from LEGEND ($10.7-22.8$ meV) [39], SNO + Phase II ($19-46$ meV) [40] and nEXO ($5.7-17.7$ meV) [41].

Note that the prediction of a lower bound for the $0\nu\beta\beta$ decay rate has been shown to occur in "missing partner" neutrino mass models, such as the "incomplete" (3,2) seesaw mechanism containing only two isosinglet neutrinos [27], or similar radiative mechanisms [28]. However here it appears in a novel way, associated with the anti-symmetry of the Yukawa coupling matrix h_{ab} determining the loop-induced neutrino mass through Eq. (12).

Note that other tree-level contributions mediated by charged scalars are neglected, since they are suppressed by the SU(3) symmetry breaking scale.

6. Flavour changing neutral currents and colliders

In this section, we comment briefly on other important phenomenological features of the present model beyond scotogenic dark matter and neutrino masses. We consider possible new contributions to well-known processes, such as meson-anti-meson mass differences, as well as genuinely new processes where the new particles present in the model could be directly produced at particle colliders such as the LHC.

In the present model flavour changing neutral current (FCNC) processes can be induced at the tree-level. Indeed, it is well-known that tree-level FCNCs appear in 331 models due to the embedding of quark families in different $SU(3)_L$ representations, as required by

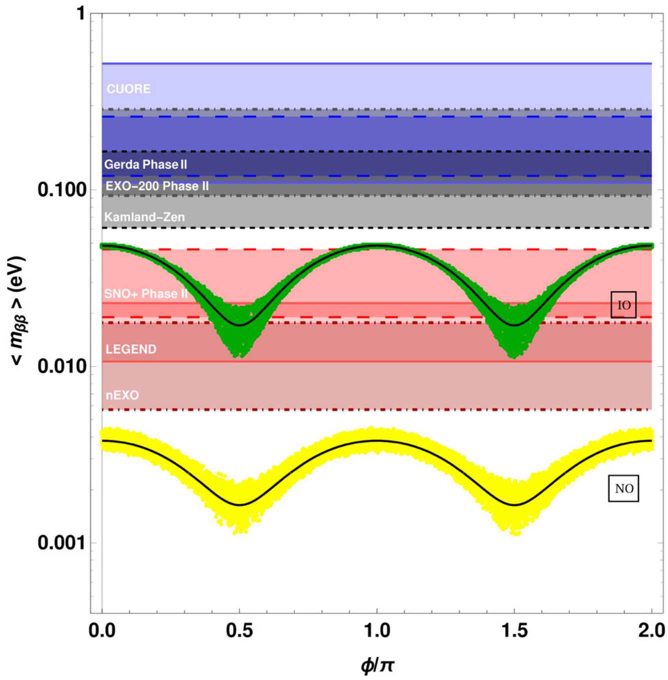

Fig. 3. Effective Majorana mass vs relative Majorana phase. The lightest neutrino is massless because of the anti-symmetry of the Yukawa coupling matrix h_{ab} determining the loop-induced neutrino mass through Eq. (12). We give the expected 3σ $0\nu\beta\beta$ bands for the case of inverted and normal mass ordering, in green and yellow, respectively. Horizontal bands represent current experimental limits and future sensitivities.

anomaly cancellation [10]. These FCNCs will be mediated by the new Z' gauge boson. In the context of these models, there have been studies devoted to such FCNCs. One finds that, in order to be in agreement with experimental results, for example $B-\bar{B}$ mass difference, the new vector bosons must be heavier than a few TeV, see Fig. 5 of Ref. [16]. Although in our extended 331 model, not one but two new vector bosons, $Z'_{1,2}$, appear [42], and both mediate FCNCs at tree-level,[3] we expect the restrictions on the intermediate vector boson masses to be similar.

Another possible origin of tree-level FCNCs is through the mixing between standard and non-standard fermions. In the present model, however, thanks to matter-parity conservation, the latter does not take place, since standard and non-standard fermions have opposite matter-parities (see Table 1) and, as such, do not mix. Therefore, matter parity conservation plays yet another role: it forbids the mixing between ordinary and new fermions, thereby preventing the appearance of potentially dangerous FCNCs associated with it.

Scalar bosons can also mediate tree-level FCNCs. In order to find the terms that lead to flavour-changing currents, we focus on the quark sector and write down the corresponding Yukawa terms

$$\begin{aligned} -\mathcal{L}_q = {} & y^d_{ia}\overline{q_{iL}}\eta^* d_{aR} + y^u_{ia}\overline{q_{iL}}\rho^* u_{aR} + y^D_{ij}\overline{q_{iL}}\chi^* D_{jR} + y^U_{i3}\overline{q_{iL}}\xi^* U_{3R} \\ & + h^u_{3a}\overline{q_{3L}}\eta u_{aR} + h^d_{3a}\overline{q_{3L}}\rho d_{aR} + h^U_{33}\overline{q_{3L}}\chi U_{3R} + h^D_{3i}\overline{q_{3L}}\xi D_{iR} + h.c.\,. \end{aligned} \tag{33}$$

The scalar singlet σ and the triplet ζ do not couple to quarks at tree level. Moreover, even though ξ couples to quarks, it does not contribute to the tree level masses as a result of matter parity conservation. Therefore, quark mass generation proceeds in a way similar to the minimal 331 models, where only three triplets are present *i.e.* η, ρ and χ.

Expanding the Yukawa operators above in terms of the field components, we find that the only neutral scalars that can mediate flavour changing currents amongst the standard quarks are η^0_1 and ρ^0_2. The relevant terms are

$$-\mathcal{L}_q \supset y^d_{ia}\overline{d_{iL}}(\eta^0_1)^* d_{aR} - y^u_{ia}\overline{u_{iL}}(\rho^0_2)^* u_{aR} + h^u_{3a}\overline{u_{3L}}\eta^0_1 u_{aR} + h^d_{3a}\overline{d_{3L}}\rho^0_2 d_{aR} + h.c.\,. \tag{34}$$

The CP-even and CP-odd components of η^0_1 and ρ^0_2 mix with other fields according to the mass matrices in Eqs. (25) and (27), respectively.

Therefore, although our model introduces new scalars when compared to the minimal 331 version, the new fields (σ, ζ and ξ) do not imply new sources of FCNCs.[4] This way, the results found for the 331 version in Ref. [43] can be adapted to our case. In that paper, the authors found that no light state at the standard model scale mediates flavour changing neutral currents. On the other hand, the heavy states, with masses at the 331 breaking scale w, do mediate FCNCs, hence consistency with experiment requires mediator masses of a few TeV or above.

We now turn to possible phenomenological implications for particle colliders. Now that Run3 of LHC is soon starting and High-Luminosity LHC is in preparation, it is very relevant to explore the possibility of producing the new particles in the currently planned experimental programme at CERN.

[3] There is also an unmixed electrically neutral gauge boson coupling neutrinos to a dark fermion.

[4] The only way some of the new scalars can contribute to tree-level FCNCs is via their mixing with η^0_1 and ρ^0_2.

First we note that the new neutral gauge bosons present in the model yield new contributions to Drell-Yan production of di-muon events at the LHC [44]. On this basis one expects bounds at the few TeV level, as seen e.g. in Fig. 5 of Ref. [16]. These are similar but complementary to the sensitivity limits obtained from meson-anti-meson mixing.

These new neutral gauge bosons also provide a portal for producing other new particles present in the model. These include the heavy quarks with electric charge 2/3 and −1/3, as well as other particles in the "dark sector". From current studies on 331 models we expect that the non-observation of any signal would restrict the parameter space, giving rise to mass limits at the few TeV level for the new particles, similar to the ones obtained above. Hence, if the masses of the new particles are chosen to be adequately large, none of these restrictions is likely to "kill the model". Instead, a number of processes associated with well-motivated TeV-scale physics could be generated within potentially achievable experimental sensitivities. In short, in addition to providing a comprehensive scotogenic framework for neutrino mass and dark matter, our model also provides a rich benchmark for new physics at collider experiments. Quantitative details require dedicated studies and simulations that lie outside the scope of this paper.

7. Summary and conclusions

Here we have proposed an $\mathrm{SU(3)_c \otimes SU(3)_L \otimes U(1)_X \otimes U(1)_N}$ electroweak extension of the standard model where dark matter stability arises from a residual matter-parity symmetry, following naturally from the spontaneous breaking of the gauge symmetry. The theory is scotogenic in the sense that dark matter is the mediator responsible for neutrino mass generation. A key feature of our new scotogenic dark matter theory is the presence of a triplet scalar boson with anti-symmetric Yukawa couplings to neutrinos. This naturally leads to a very simple characteristic prediction, i.e. one of the light neutrinos is massless, thus implying a lower bound for the $0\nu\beta\beta$ decay rate. In contrast to most other models where a massless neutrino arises from an *ad hoc* incomplete multiplet choice, here it is an unavoidable characteristic feature of the theory. The theory also provides a comprehensive framework for scotogenic dark matter and a rich benchmark for FCNC tests and collider searches at the LHC.

Acknowledgements

Work supported by the Spanish grants SEV-2014-0398 and FPA2017-85216-P (AEI/FEDER, UE), PROMETEO/2018/165 (Generalitat Valenciana) and the Spanish Red Consolider MultiDark FPA2017-90566-REDC. J. L. acknowledges financial support under grant 2019/04195-7, São Paulo Research Foundation (FAPESP), while OP is supported by the National Research Foundation of Korea, under Grants No. 2017K1A3A7A09016430 and No. 2017R1A2B4006338. The Feynman diagram is drawn using TikZ-Feynman [45].

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